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# Electroweak Chiral Lagrangians and $\gamma\gamma \rightarrow Z_L Z_L, W_L^+ W_L^-$ at One Loop

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## ABSTRACT

In these proceedings we discuss our recent work on  $\gamma\gamma \rightarrow W_L^+ W_L^-$  and  $\gamma\gamma \rightarrow Z_L Z_L$  within the framework of Electroweak Chiral Lagrangians with a light Higgs. These observables are good candidates to provide indications of physics beyond the Standard Model and can complement other analyses and global fits. Making use of the equivalence theorem, we have performed the computation up to the next-to-leading order in the chiral expansion, i.e., including one-loop corrections and the full set of possible counterterms allowed at that order in the low-energy effective field theory. The one-loop amplitudes turn out to have a extremely simple structure and they are ultraviolet finite. Thus, the relevant combinations of higher order chiral couplings  $c_\gamma^r$  and  $(a_1^r - a_2^r + a_3^r)$  do not run.

We also identify an additional set of observables that depend on these same parameters, which can be use to complement the  $\gamma\gamma$  analysis through a global fit.

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# 1 Introduction

Based on the article [1], we discuss in these proceedings our study of the photon-photon scattering into longitudinal weak gauge bosons,  $\gamma\gamma \rightarrow Z_L Z_L, W_L^+ W_L^-$ . This observable can be an optimal place to look for deviations from the Standard Model (SM) in future analyses and to test the composite nature of the Higgs. Since this boson does not contribute at tree-level to the photon-photon scattering in the SM, any new physics in the Higgs is expected to give particular signatures. In combination with other observables it is possible to extract the relevant couplings of the low-energy effective field theory (EFT), the Electroweak Chiral Lagrangian including the Higgs (ECLh) [1].

On the experimental side, there are dedicated programs at LHC on photon-photon scattering. The CMS Collaboration has published the first results on  $\gamma\gamma \rightarrow W^+ W^-$  [4], showing the feasibility of this type of analysis. There are indeed the approved projects CMS-TOTEM [2] and ATLAS-AFP [3], where new forward proton detectors will be incorporated (if funding is finally approved). Although so far the statistics are very low these are expected to be increased and become more relevant in future runs, as one is able to reach higher energies [5].

So far, LHC data have clearly indicated the presence of a relatively light Higgs-like scalar ( $m_h \simeq 126$  GeV) with coupling strengths close to the SM ones. E.g. one has  $a \approx 1$  for the  $hWW$  coupling  $a$  ( $\kappa_V$  in other notations [6, 7]), within  $\mathcal{O}(10\%)$  uncertainties:  $a = 1.15 \pm 0.08$  (ATLAS [6]),  $a = 1.01 \pm 0.07$  (CMS [7]). Notice that a deviation from its SM value  $a = 1$  implies immediately a bad ultraviolet behaviour of the electroweak (EW) boson scattering amplitudes, which needs of new beyond-SM (BSM) states to restore unitarity. Thus, theoretical analyses have been performed on the  $WW$ ,  $ZZ$  and  $hh$  scattering amplitudes [8] and ATLAS has recently provided the evidence of  $W^\pm W^\pm$  scattering and the first rough bounds on the relevant low-energy ECLh parameters [9].

In order to pin down these subtle effects one needs to go beyond using simple effective vertices to describe transitions such as  $\gamma\gamma \rightarrow W^+ W^-$ . Instead one needs to properly incorporate the EW loops (both UV-divergences and finite parts), being the ECLh framework the most appropriate one in the current situation where there seems to be a large mass gap and no new BSM particle has been found below the TeV. This also means that in general one has to keep in the phenomenological analysis all the counterterms allowed by the symmetry, unless one considers particular BSM models.

For the construction of the low-energy EFT Lagrangian we will consider just the observed particles, i.e., the SM content. More precisely, we will restrict our calculation to the bosonic sector including the Higgs and the EW gauge bosons  $W^a$  and Would-Be-Goldstone Bosons (WBGB)  $w^a$ . We will base our EFT on a custodial symmetry pattern, where the spontaneous symmetry breaking (SSB)  $SU(2)_L \times SU(2)_R / SU(2)_{L+R}$  gives place to the WBGB. We will not specify the origin of the SSB and the WBGB will be non-linearly realized. For simplicity, we perform our one-loop analysis in the Landau gauge, where the ghost sector decouples and the WBGBs become massless. In any case, in order to have a well defined classification of the scaling of the diagrams one must work within renormalizable  $R_\xi$  gauges, which allow us to use the Equivalence Theorem (Eq.Th.) [11] and ensure an adequate  $1/k^2$  scaling of the EW gauge boson propagators in the UV [1, 10]. Thus, we will make use of the Eq.Th. [11],

$$\mathcal{M}(\gamma\gamma \rightarrow W_L^a W_L^b) \simeq -\mathcal{M}(\gamma\gamma \rightarrow w^a w^b), \quad (1)$$

and extract the observable in the energy regime  $m_W^2, m_Z^2 \ll s$ , with  $W^a$  and  $w^a$  the EW gauge bosons and the WBGB, respectively. The EW gauge boson masses  $m_{W,Z}$  are then neglected in our computation. \* Furthermore, since  $m_h \sim m_{W,Z} \ll 4\pi v = 3$  TeV (with the vacuum expectation value  $v = 246$  GeV) we also neglect the Higgs mass in our calculation, which means not including Higgs mass corrections, typically proportional to  $\mathcal{O}(m_h^2/(16\pi^2 v^2))$  [1, 8]. In summary, the applicability range of our analysis [1] is

$$m_W^2, m_Z^2, m_h^2 \stackrel{\text{Eq.Th.}}{\ll} s, t, u \stackrel{\text{EFT}}{\ll} \Lambda_{\text{ECLh}}^2, \quad (2)$$

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\* One must be careful with this limit, since  $m_W^2 = g^2 v^2/4$ ,  $m_Z^2 = (g^2 + g'^2)v^2/4$  are set to zero while  $v$  is kept fixed. This means we are taking  $g, g' \rightarrow 0$  and considering just the leading non-vanishing contribution to the photon scattering amplitude, neglecting higher powers of  $g^{(\prime)}$ .

with the upper limit given by the EFT cut-off  $\Lambda_{\text{ECLh}}$ , expected to be of the order of  $4\pi v = 3 \text{ TeV}$  or the mass of possible heavy BSM particles. In principle, a more complete calculation beyond Eq.Th. would allow to extend our ECLh result up to near-threshold energies where, however, BSM effects are much more suppressed.

## 2 The ECLh Lagrangian relevant for $\gamma\gamma \rightarrow w^a w^b$

The ECLh Lagrangian contains as dynamical fields the EW gauge bosons,  $W^\pm$ ,  $Z$  and  $\gamma$ , the corresponding WBGBs ( $w^\pm$  and  $z$ ) and the Higgs-like scalar boson,  $h$ . The WBGBs are described by a matrix field  $U$  that takes values in the  $SU(2)_L \times SU(2)_R / SU(2)_{L+R}$  coset, and transforms as  $U \rightarrow LUR^\dagger$  under the action of the global group  $SU(2)_L \times SU(2)_R$  that defines the EW Chiral symmetry [12, 13]. The relevant ECLh Lagrangian with the basic building blocks are

$$\begin{aligned} D_\mu U &= \partial_\mu U + i\hat{W}_\mu U - iU\hat{B}_\mu, & V_\mu &= (D_\mu U)U^\dagger, \\ \hat{W}_{\mu\nu} &= \partial_\mu \hat{W}_\nu - \partial_\nu \hat{W}_\mu + i[\hat{W}_\mu, \hat{W}_\nu], & \hat{B}_{\mu\nu} &= \partial_\mu \hat{B}_\nu - \partial_\nu \hat{B}_\mu, & \hat{W}_\mu &= gW_\mu^a \tau^a/2, & \hat{B}_\mu &= g' B_\mu \tau^3/2. \end{aligned} \quad (3)$$

Two particular parametrizations of the unitary matrix  $U$  in terms of the dimensionless  $w^\pm/v$  and  $z/v$  fields (exponential, with  $U = \exp\{i\tau^a w^a/v\}$ , and spherical, with  $U = \sqrt{1 - w^a w^a/v^2} + i\tau^a w^a/v$ ), were considered in [1], both leading to the same predictions for the physical (on-shell) observables.<sup>†</sup> In order to have a power counting consistent with the loop expansion we consider  $\partial_\mu, (gv), (g'v), m_h \sim \mathcal{O}(p)$  or, equivalently,  $\partial_\mu, m_W, m_Z, m_h \sim \mathcal{O}(p)$ , and the scaling of the tensors  $D_\mu U, V_\mu \sim \mathcal{O}(p)$  and  $\hat{W}_{\mu\nu}, \hat{B}_{\mu\nu} \sim \mathcal{O}(p^2)$  [1, 10].

With these building blocks one then constructs the ECLh up to a given order in the chiral expansion. We require this Lagrangian to be CP invariant, Lorentz invariant and  $SU(2)_L \times U(1)_Y$  gauge invariant. We focus ourselves on the relevant terms for  $\gamma\gamma \rightarrow w^+ w^-$  and  $\gamma\gamma \rightarrow zz$  scattering processes. First, the relevant terms in the leading order (LO) Lagrangian – of  $\mathcal{O}(p^2)$  – are given by

$$\mathcal{L}_2 = -\frac{1}{2g^2} \text{Tr}(\hat{W}_{\mu\nu} \hat{W}^{\mu\nu}) - \frac{1}{2g'^2} \text{Tr}(\hat{B}_{\mu\nu} \hat{B}^{\mu\nu}) + \frac{v^2}{4} \left[ 1 + 2a \frac{h}{v} + b \frac{h^2}{v^2} \right] \text{Tr}(D^\mu U^\dagger D_\mu U) + \frac{1}{2} \partial^\mu h \partial_\mu h + \dots, \quad (5)$$

where the dots stand for  $\mathcal{O}(p^2)$  operators with three or more Higgs fields, which do not enter into this computation. As already said, the Higgs mass term is neglected in our analysis [1]. At NLO, we will have contributions from the  $\mathcal{O}(p^4)$  Lagrangian [1, 13],

$$\mathcal{L}_4 = a_1 \text{Tr}(U \hat{B}_{\mu\nu} U^\dagger \hat{W}^{\mu\nu}) + ia_2 \text{Tr}(U \hat{B}_{\mu\nu} U^\dagger [V^\mu, V^\nu]) - ia_3 \text{Tr}(\hat{W}_{\mu\nu} [V^\mu, V^\nu]) - \frac{c_\gamma}{2} \frac{h}{v} e^2 A_{\mu\nu} A^{\mu\nu} + \dots \quad (6)$$

involving the photon field strength,  $A_{\mu\nu} = \partial_\mu A_\nu - \partial_\nu A_\mu$ . The dots stand for other  $\mathcal{O}(p^4)$  operators not contributing to  $\gamma\gamma$  scattering. Notice that the singlet Higgs field  $h$  enters in the ECLh Lagrangian via multiplicative polynomial functions and their derivatives.

## 3 Analytical results for $\gamma\gamma \rightarrow w^a w^b$ in ECLh up to NLO

In dimensional regularization, our NLO computation of the  $\mathcal{M}(\gamma\gamma \rightarrow w^a w^b)$  amplitudes can be systematically sorted out in the form [1]

$$\mathcal{M} = \mathcal{M}_{\text{LO}} + \mathcal{M}_{\text{NLO}} \sim \underbrace{\mathcal{O}(e^2)}_{\text{LO, tree}} + \left( \underbrace{\mathcal{O}\left(e^2 \frac{p^2}{16\pi^2 v^2}\right)}_{\text{NLO, 1-loop}} + \underbrace{\mathcal{O}\left(e^2 \frac{a_i p^2}{v^2}\right)}_{\text{NLO, tree}} \right). \quad (7)$$

The LO contributions are given by all the tree-level diagrams with vertices from just the  $\mathcal{L}_2$  Lagrangian. The NLO contributions come from two types of graphs: one-loop diagrams with only vertices from  $\mathcal{L}_2$  (which

<sup>†</sup> Other representations have been recently studied in Ref. [14].

may *a priori* generate UV divergences); tree-level diagrams with only one vertex from the  $\mathcal{L}_4$  Lagrangian, being the remaining ones from  $\mathcal{L}_2$ . The  $a_i$  represent generic  $\mathcal{L}_4$  couplings ( $a_1, a_2, a_3, c_\gamma$  in our case) and will renormalize the referred one-loop UV divergences through an appropriate renormalization  $a_i^r = a_i + \delta a_i$  [1, 8, 14], following a procedure completely analogous to that in Chiral Perturbation Theory [15]. The  $\delta a_i$  counterterms in turn, will lead to the corresponding running of the renormalized ECLh parameters, i.e. the dependence with the renormalization scale  $\mu$  of the  $a_i^r(\mu)$ .

The amplitudes  $\mathcal{M}(\gamma(k_1, \epsilon_1)\gamma(k_2, \epsilon_2) \rightarrow w^a(p_1)w^b(p_2))$  with  $w^a w^b = zz, w^+ w^-$ , are given in terms of the Lorentz decomposition [1]

$$\mathcal{M} = ie^2(\epsilon_1^\mu \epsilon_2^\nu T_{\mu\nu}^{(1)})A(s, t, u) + ie^2(\epsilon_1^\mu \epsilon_2^\nu T_{\mu\nu}^{(2)})B(s, t, u), \quad (8)$$

written in terms of the two independent Lorentz structures involving the external momenta, which can be found in [1]. The Mandelstam variables are defined as  $s = (p_1 + p_2)^2$ ,  $t = (k_1 - p_1)^2$  and  $u = (k_1 - p_2)^2$  and the  $\epsilon_i$ 's are the polarization vectors of the external photons.

For the scalar functions  $A$  and  $B$  we will follow the same notation as in Eq. (7):  $A = A_{\text{LO}} + A_{\text{NLO}}$ ,  $B = B_{\text{LO}} + B_{\text{NLO}}$ . Thus, At LO one has the well-known tree-level contribution from  $\mathcal{L}_2$ ,

$$A(\gamma\gamma \rightarrow zz)_{\text{LO}} = B(\gamma\gamma \rightarrow zz)_{\text{LO}} = 0, \quad (9)$$

$$A(\gamma\gamma \rightarrow w^+ w^-)_{\text{LO}} = 2sB(\gamma\gamma \rightarrow w^+ w^-)_{\text{LO}} = -\frac{1}{t} - \frac{1}{u}. \quad (10)$$

At NLO, i.e at  $\mathcal{O}(e^2 p^2)$ , we find

$$A(\gamma\gamma \rightarrow zz)_{\text{NLO}} = \frac{2ac_\gamma^r}{v^2} + \frac{(a^2 - 1)}{4\pi^2 v^2}, \quad B(\gamma\gamma \rightarrow zz)_{\text{NLO}} = 0, \quad (11)$$

$$A(\gamma\gamma \rightarrow w^+ w^-)_{\text{NLO}} = \frac{2ac_\gamma^r}{v^2} + \frac{(a^2 - 1)}{8\pi^2 v^2} + \frac{8(a_1^r - a_2^r + a_3^r)}{v^2}, \quad B(\gamma\gamma \rightarrow w^+ w^-)_{\text{NLO}} = 0. \quad (12)$$

The term with  $c_\gamma^r$  comes from the Higgs tree-level exchange in the  $s$ -channel, the term proportional to  $(a^2 - 1)$  comes from the one-loop diagrams with  $\mathcal{L}_2$  vertices, and the Higgsless operators in Eq. (6) yield the tree-level contribution to  $\gamma\gamma \rightarrow w^+ w^-$  proportional to  $(a_1 - a_2 + a_3)$ . It is worthy to remark that the one-loop amplitude does not carry the typical chiral suppression  $\mathcal{O}(E^2/(16\pi^2 v^2))$  [16] (with  $E^2 = s, t, u$ ) but a stronger one  $\mathcal{O}((a^2 - 1)E^2/(16\pi^2 v^2))$ . Notice that the experimental determinations of  $a$  find it to be close to 1 within an  $\mathcal{O}(10\%)$  uncertainty [6, 7]. Thus, in the limit  $a \rightarrow 1$  the loop contribution vanishes, as one finds in the SM under the same approximations consider in our work [1].

Independent diagrams are in general UV divergent and have complicated Lorentz structures and energy dependence (see App. B in [1] for details). However, in dimensional regularization, the final one-loop amplitude turns out to be UV finite after all the contributions are put together. The reason for this subtle cancellation can be understood from the equivalence between our loop diagrams and the one-loop amplitude in the  $SO(5)/SO(4)$  Non-Linear Sigma Model [1, 18] in the approximations considered in our analysis [1]. Thus, the combinations of  $\mathcal{L}_4$  chiral parameters  $(a_1^r - a_2^r + a_3^r)$  and  $c_\gamma^r$  do not need to be renormalized:

$$a_1^r - a_2^r + a_3^r = a_1 - a_2 + a_3, \quad c_\gamma^r = c_\gamma. \quad (13)$$

Notice that when setting  $a = c_\gamma = 0$  in our formulas above we recover exactly the previous result found in [16] for the case of the Higgsless EW Chiral Lagrangian (ECL) and the analogous result for the amplitude in the pion case [17].

## 4 Related observables, global analysis and running determination

As we can see in Eq. (12) the  $\gamma\gamma \rightarrow zz$  and  $\gamma\gamma \rightarrow w^+ w^-$  cross sections depend on three independent combinations of parameters:  $a$ ,  $c_\gamma$  and  $(a_1 - a_2 + a_3)$ . In order to pin down each of these ECLh couplings one must combine our  $\gamma\gamma$ -scattering analysis with other observables that depend on this same set of parameters. It is not difficult to find that other processes involving photons depend on these parameters. In Ref. [1]

Observables	Relevant combinations of parameters	
	from $\mathcal{L}_2$	from $\mathcal{L}_4$
$\mathcal{M}(\gamma\gamma \rightarrow zz)$	$a$	$c_\gamma^r$
$\mathcal{M}(\gamma\gamma \rightarrow w^+w^-)$	$a$	$(a_1^r - a_2^r + a_3^r), c_\gamma^r$
$\Gamma(h \rightarrow \gamma\gamma)$	$a$	$c_\gamma^r$
$S$ -parameter	$a$	$a_1^r$
$\mathcal{F}_{\gamma^*ww}$	$a$	$(a_2^r - a_3^r)$
$\mathcal{F}_{\gamma^*\gamma h}$	$-$	$c_\gamma^r$

Table 1: Set of six observables studied in Ref. [1] and their corresponding relevant combinations of chiral parameters.

	ECLh	ECL (Higgsless)
$\Gamma_{a_1 - a_2 + a_3}$	0	0
$\Gamma_{c_\gamma}$	0	-
$\Gamma_{a_1}$	$-\frac{1}{6}(1 - a^2)$	$-\frac{1}{6}$
$\Gamma_{a_2 - a_3}$	$-\frac{1}{6}(1 - a^2)$	$-\frac{1}{6}$
$\Gamma_{a_4}$	$\frac{1}{6}(1 - a^2)^2$	$\frac{1}{6}$
$\Gamma_{a_5}$	$\frac{1}{8}(b - a^2)^2 + \frac{1}{12}(1 - a^2)^2$	$\frac{1}{12}$

Table 2: Running of the relevant ECLh parameters and their combinations appearing in the six selected observables. The third column provides the corresponding running for the Higgsless ECL case [19].

we computed 4 more observables of this kind: the  $h \rightarrow \gamma\gamma$  decay width, the oblique  $S$ -parameter, and the  $\gamma^* \rightarrow w^+w^-$  and  $\gamma^*\gamma \rightarrow h$  electromagnetic form-factors. In Table 1 we provide the relevant combinations of  $\mathcal{L}_2$  and  $\mathcal{L}_4$  coefficients in each case. The detailed amplitudes can be found in App. D in [1]. These six observables allow the extraction of the four independent combinations of couplings  $a, c_\gamma, a_1$  and  $(a_2 - a_3)$ . Moreover, one can use four of these observables to perform a prediction for the remaining two.

The one-loop contribution in the six relevant amplitudes is found to be UV-divergent in some cases. These divergences are absorbed by means of the renormalizations  $a_i^r(\mu) = a_i + \delta a_i$  and the  $\mathcal{O}(p^4)$  counterterms  $\delta a_i$  (with  $a_i$  referring to the corresponding  $\mathcal{L}_4$  coefficients  $a_i = c_\gamma, a_1, a_2, a_3$ ). As expected, the renormalization in the six observables gives a fully consistent set of renormalization conditions and fixes the running of the renormalized couplings in the way

$$\frac{da_i^r}{d\ln\mu} = -\frac{\Gamma_{a_i}}{16\pi^2}, \quad (14)$$

with the corresponding  $\Gamma_{a_i}$  given in Table 2. For the sake of completeness, we have added the running of the ECLh parameters  $a_4^r$  and  $a_5^r$ , which have been recently determined in the one-loop analysis of  $WW$ -scattering within the framework of chiral Lagrangians [8]. One can see that in the SM limit  $a = b = 1$  the  $\mathcal{L}_4$  coefficients in Table 2 do not run, in agreement with the fact that these higher order operators are absent in the SM.

To end up the discussion we would like to point out a series of improvements that will be incorporated to the  $\gamma\gamma$  analysis [1] in a future phenomenological analysis. First, one must take into account fermion loops and in particular the top quark contribution, which is expected to be as numerically important as the  $W$  loops, as it happens in analogous processes like e.g.  $h \rightarrow \gamma\gamma$ . Likewise, the computation can be refined by taking into account the impact of the  $W, Z$  and  $h$  masses, allowing the extension of the analysis all the way down to the production threshold. Thus, one should go beyond the Equivalence Theorem and compute the actual  $\gamma\gamma \rightarrow W_L^a W_L^b$  amplitudes. These light mass effects are nevertheless suppressed below the percent level for energies of the order of the TeV.

We plan to perform a MonteCarlo analysis for LHC and future  $e^+e^-$  accelerators in order to estimate both integrated even rates and the energy dependence. Previous studies expect a larger relevance of the  $\gamma\gamma \rightarrow W^+W^-$  subprocess in the  $W^+W^-$  production at future LHC runs, which can even exceed the  $gg \rightarrow W^+W^-$  contribution to the cross section at large  $p_T$  [5].

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